

# Curvaton reheating allows TeV Hubble scale in NO inflation

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(Dated: February 1, 2008)

Curvaton reheating is studied in non-oscillatory (NO) models of inflation, with the aim to obtain bounds on the parameters of curvaton models and find out whether low scale inflation can be attained. Using a minimal curvaton model, it is found that the allowed parameter space is considerably larger than in the case of the usual oscillatory inflation models. In particular, inflation with Hubble scale as low as 1 TeV is comfortably allowed.

Inflation is to date the most compelling solution for the horizon and flatness problems of big bang cosmology. Recently, observational data at high precision have confirmed that structure formation and the CMB anisotropy are due to the existence of superhorizon curvature perturbations, which need an acausal mechanism to be created and, therefore, are further evidence of inflation.

The majority of inflationary models suggest that inflation took place at energy comparable to that of grand unification. This is because, in most inflationary models, the curvature perturbations are due to the quantum fluctuations of the inflaton field; the field controlling the dynamics of inflation. In this case the inflationary energy scale is [1]

$$V_*^{1/4} = 0.027 \epsilon_*^{1/4} m_P, \quad (1)$$

where  $m_P = 2.4 \times 10^{18}$  GeV is the reduced Planck mass and  $\epsilon \equiv -\dot{H}/H^2$  is one of the so-called slow-roll parameters, where  $H \equiv \dot{a}/a$  with  $a$  being the scale factor and the dot denoting derivative with respect to time. In the above, ‘\*’ denotes the epoch when the cosmological scales exit the horizon during inflation. For inflation to occur one needs  $\epsilon < 1$ , which is equivalent to  $\ddot{a} > 0$ . In slow-roll models of inflation  $\epsilon \ll 1$ , which means  $H \simeq \text{constant}$ . However, slow-roll inflation is not exactly de Sitter expansion, because it has to end for the hot big bang to begin. Therefore, albeit small, in most models  $\epsilon$  is not tiny. In fact, it turns out that, typically,  $\epsilon_* \sim 1/N_*$ , where  $N_*$  is the remaining number of inflation e-folds when the cosmological scales exit the horizon.  $N_*$  ranges between 40-65 for inflation scales between grand unification and Big Bang Nucleosynthesis (BBN). Hence, from Eq. (1) we see that  $V_*^{1/4} \sim 10^{15-16}$  GeV. Exceptions to this rule exist, where  $\epsilon$  is exponentially suppressed due to appropriate features of the scalar potential (e.g. see Ref. [2]) but they are typically fine-tuned (e.g. see Ref. [3]).

However, in recent years, advances in string theory have turned the attention at progressively lower energy scales. In particular, large extra dimensions suggest that the fundamental scale may be much smaller than  $m_P$

and the scale of grand unification, possibly near the electroweak scale instead (e.g. see [4]). This means that the hot big bang cannot be extended up to grand unification energies and so inflation must take place at lower scale. Moreover, in the search for the inflaton field, the most promising candidates are either flat directions in supersymmetric theories or string moduli fields, which correspond to the size and shape of the compactified extra dimensions. However, in the MSSM and its extensions the corresponding scales are typically much less than the grand unified scale. Similarly, for string moduli one expects the scalar potential to vary significantly (i.e.  $\delta V/V \sim 1$ ) over distances  $\sim m_P$  in field space, which means that the typical scale for modular inflation is  $V_*^{1/4} \sim \sqrt{m_{3/2} m_P} \sim 10^{10.5}$  GeV (i.e.  $H_* \sim 1$  TeV) for moduli with masses of order  $m_{3/2} \sim 1$  TeV, such as string axions. Recently, many possibilities for inflation in string theory have arisen in the context of the string landscape [5]. For example, metastable supersymmetric vacua [6] may account for the vacuum density of inflation. This sets the value of  $V_*$  at the supersymmetry breaking scale which corresponds to the above mentioned scale or lower. This is why there has been growing interest in achieving inflation at low energy scales.

One way to liberate inflation from the constraint in Eq. (1) is to consider that the curvature perturbations generated by inflation are due to the quantum fluctuations of a field *other* than the inflaton, in which case Eq. (1) turns into an upper bound [7]. This, so called curvaton field, does not influence the dynamics of inflation but becomes important after inflation has ended, when it imprints its curvature perturbation onto the Universe [8]. Under this hypothesis, it is possible to relax this constraint substantially when the curvaton model is specially designed to do so [9, 10]. However, generically, for a minimal curvaton model, there is still a lower bound on the inflationary scale of order  $V_*^{1/4} \gtrsim 10^{12}$  GeV (i.e.  $H_* \gtrsim 10^6$  GeV) [11], when considering inflation, which ends through oscillations of the inflaton field, as usually assumed.

In this letter we show that, for a minimal curvaton model, the inflation scale can be substantially lower when considering a non-oscillatory (NO) model of inflation. In such a model the inflaton corresponds to a flat direction with runaway behaviour, typical, for example, of the

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scalar potential of the so-called geometric string moduli fields. In such models after inflation ends the Universe remains dominated by the inflaton field which is now fast-rolling down its runaway scalar potential [12]. Since the inflaton does not oscillate and does not decay after the end of inflation, reheating occurs by other means.<sup>1</sup> Recently, a curvaton field has been assumed to be responsible for reheating in NO inflation models [13, 14, 15]. However, up to now the parameter space for the inflationary scale (in particular low-scale inflation) has not been studied.

We begin by outlining the particulars of NO inflation. In NO models, after the end of inflation, the inflaton field becomes dominated by its kinetic density  $\rho_\phi \sim \rho_{\text{kin}} \equiv \frac{1}{2}\dot{\phi}^2$  [12]. The equation of motion for the field becomes  $\ddot{\phi} + 3H\dot{\phi} \simeq 0$ , which is oblivious to the scalar potential and suggests that  $\rho_\phi \propto a^{-6}$ . This gives rise to a period of so-called kination, where  $a \propto t^{1/3}$  [16]. Because the dilution of the inflaton's density is so drastic, the Universe can be reheated by the thermal bath due to the decay products of a curvaton field, which also accounts for the density perturbations in the Universe [15]. We should stress here that our treatment below is independent of the NO inflation model since the scalar potential for the inflaton sector is negligible during the evolution of the curvaton field.

We consider a minimal curvaton scenario where the curvaton field  $\sigma$  is characterised by the scalar potential

$$V(\sigma) = \frac{1}{2}m^2\sigma^2. \quad (2)$$

In order to act as a curvaton  $\sigma$  must undergo particle production during inflation. Hence, it needs to be effectively massless. This implies the bound  $m < H_*$ . An even tighter bound, however, is due to spectral index considerations. Indeed, for the curvaton we have  $n_s - 1 \simeq 2\eta = \frac{2}{3}(m/H_*)^2$ , where  $\eta \equiv m_P^2 \frac{1}{V_*} \frac{\partial^2 V}{\partial \sigma^2}$  [8] with  $V_* \simeq 3H_*^2 m_P^2$  being the density during inflation. Since observations do not favour a blue spectrum, we require that  $n_s$  is at most  $n_s \simeq 1.00$ , which is still marginally acceptable.<sup>2</sup> This means that  $m$  must satisfy the bound

$$m \leq 0.1 H_* . \quad (3)$$

We assume that the inflaton's contribution to the curvature perturbation is negligible, as is the norm under the curvaton hypothesis [8]. In any case, it is difficult to obtain a sizable curvature perturbation in low scale inflation [7] (See, however, Ref. [2]). Since the Universe is reheated by the decay products of the curvaton which

dominate the scalar field background, the observed curvature perturbation  $\zeta = 4.8 \times 10^{-5}$  is entirely due to the curvaton, i.e.

$$\begin{aligned} \zeta &= \zeta_\sigma \equiv -H \left. \frac{\delta\rho_\sigma}{\dot{\rho}_\sigma} \right|_{\text{dec}} = \frac{1}{3} \left. \frac{\delta\rho_\sigma}{\rho_\sigma} \right|_{\text{dec}} = \frac{2}{3} \left. \frac{\delta\sigma}{\sigma} \right|_{\text{dec}} \\ &= \frac{2}{3} \left. \frac{\delta\sigma}{\sigma} \right|_{\text{osc}} \simeq \frac{2}{3} \left. \frac{\delta\sigma}{\sigma} \right|_* = \frac{H_*}{3\pi\sigma_*} \Rightarrow \sigma_* \sim \frac{H_*}{\zeta}, \end{aligned} \quad (4)$$

where we used that, before its decay (denoted as 'dec'), the curvaton undergoes quasi-harmonic oscillations in the scalar potential of Eq. (2), so that  $\rho_\sigma = 2\overline{V(\sigma)} \propto a^{-3}$  [17], where  $\overline{V(\sigma)}$  is the average value of the potential density over many oscillations. Since the oscillating curvaton satisfies the same linear equation of motion as its oscillating perturbation, (namely  $\ddot{\sigma} + 3H\dot{\sigma} + m^2\sigma = 0$ ) the fractional perturbation  $\delta\sigma/\sigma$  remains constant and equal to its value at the onset of the oscillations, denoted as 'osc'. Before the onset of the oscillations, both the value of the field and its perturbation are frozen so the fractional perturbation remains constant again and equal to its value at horizon exit, denoted by '\*'. As is the case with particle production of scalar fields, the value of the curvaton perturbation at horizon exit is determined by the Hawking temperature:  $\delta\sigma_* = H_*/2\pi$ . Note that Eq. (4) implies that  $\sigma_* \gg H_*$ , which guarantees that the curvature perturbations are predominantly Gaussian [8].

After the end of inflation, the inflaton becomes kinetically dominated, while the curvaton remains frozen because  $m \ll H_*$  and its motion is overdamped. This means that the density parameter of the curvaton  $\Omega \equiv \rho_\sigma/\rho$  scales as  $\Omega \propto a^6$  until  $H(t)$  drops down to  $H \sim m$ , when the curvaton unfreezes and begins its quasi-harmonic oscillations. At this time we have

$$\Omega_{\text{osc}} \sim \left( \frac{\sigma_*}{m_P} \right)^2. \quad (5)$$

After the onset of the oscillations, the curvaton density decreases as  $\rho_\sigma \propto a^{-3}$  so that  $\Omega \propto a^3$ .

Now, there are two possibilities, depending on whether the curvaton decays before its density dominates the Universe or afterwards. Let us assume firstly that the curvaton decays before domination. In this case, using that during kination  $a \propto H^{-1/3}$ , one readily obtains that the curvaton density parameter at the time of curvaton decay is

$$\Omega_{\text{dec}} \sim \frac{m}{\Gamma} \left( \frac{\sigma_*}{m_P} \right)^2, \quad (6)$$

where  $\Gamma$  is the curvaton decay rate. The decay products of the curvaton redshift as relativistic matter  $\rho_\gamma \propto a^{-4}$ , which means that their density parameter grows as  $\Omega_\gamma \propto a^2$ . Reheating is achieved when this density parameter reaches unity and the Universe becomes dominated by the thermal bath of the curvaton decay products. Using that  $H_{\text{dom}} \sim T_{\text{reh}}^2/m_P$ , it is easy to find

$$T_{\text{reh}} \sim \sqrt{m_P \Gamma} \left( \frac{m}{\Gamma} \right)^{3/4} \left( \frac{H_*}{\zeta m_P} \right)^{3/2}, \quad (7)$$

<sup>1</sup> NO models have been particularly useful for quintessential inflation (e.g. see Ref. [13, 14] and references therein) because the inflaton field can survive until today and be responsible for the dark energy at present.

<sup>2</sup> Smaller values of  $n_s$  can be attained if  $\epsilon$  is not tiny, in which case  $n_s - 1 \simeq 2(\eta - \epsilon)$  [8].

where we also employed Eq. (4) and we denoted with ‘dom’ the epoch when the inflaton’s density becomes subdominant. From the above and using also that  $\Gamma \geq H_{\text{dom}}$  and  $T_{\text{reh}} > T_{\text{BBN}}$  we obtain the bound

$$H_* > (10\zeta^2)^{1/3} (T_{\text{BBN}}^2 m_P)^{1/3} \sim 10 \text{ GeV}, \quad (8)$$

where we also considered Eq. (3) and demanded that reheating occurs before BBN; the latter corresponding to temperature  $T_{\text{BBN}} \sim 1 \text{ MeV}$ . Thus, we see that inflation with  $H_* \sim 1 \text{ TeV}$  is possible to accommodate with a minimal curvaton model.

Let us consider, now, the case when the curvaton decays after it dominates the Universe. In this case the oscillating curvaton dominates when

$$H_{\text{dom}} \sim m \left( \frac{\sigma_*}{m_P} \right)^2. \quad (9)$$

Reheating is, now, achieved at the decay of the curvaton, so that

$$T_{\text{reh}} \sim \sqrt{m_P \Gamma}. \quad (10)$$

Using that  $\Gamma \leq H_{\text{dom}}$  and also  $T_{\text{reh}} > T_{\text{BBN}}$ , we arrive once more at the bound in Eq. (8), where we also considered Eq. (3).

Additional bounds on the parameters are obtained as follows. For the decay rate of the curvaton we can write  $\Gamma = g^2 m$ , where  $g$  is the coupling of the curvaton to its decay products. The expected range for this coupling is:

$$\frac{m}{m_P} \lesssim g \lesssim 1, \quad (11)$$

where the lower bound is due to gravitational decay.

Assume at first that the curvaton decays before domination. Then, from Eq. (7), we have  $T_{\text{reh}} \propto \Gamma^{-1/4}$ . Consequently, combining the lower bound in Eq. (11) with the requirement that  $T_{\text{reh}} > T_{\text{BBN}}$  and Eq. (7), we obtain the bound

$$H_* \gtrsim \zeta (T_{\text{BBN}}^2 m_P)^{1/3} \sim 1 \text{ GeV}, \quad (12)$$

which is weaker than the bound in Eq. (8). Also, using that  $T_{\text{reh}} > T_{\text{BBN}}$ , Eq. (7) results in

$$g < \frac{m_P m}{T_{\text{BBN}}^2} \left( \frac{H_*}{\zeta m_P} \right)^3 \leq 10^{41} \zeta^{-3} \left( \frac{H_*}{m_P} \right)^4, \quad (13)$$

where, in the last inequality, we used Eq. (3). Comparing this bound with the lower bound in Eq. (11) it can be easily verified (with the help of Eq. (12)) that there is always parameter space for  $g$ . Using also that  $g \lesssim 1 \leq 0.1 H_*/m$  according to Eqs. (3) and (11), it can be shown that the bound in Eq. (13) is tighter than the upper bound in Eq. (11) only for  $H_* < 10 \text{ TeV}$ . Finally, since in the case of decay before domination we have  $\Gamma \geq H_{\text{dom}} \sim T_{\text{reh}}^2/m_P$ , employing Eq. (7) the upper bound in Eq. (11) suggests

$$\sigma_* \lesssim m_P \Rightarrow H_* \lesssim \zeta m_P \sim 10^{14} \text{ GeV}, \quad (14)$$

where we also used Eq. (4).

Now, let us consider the case when the curvaton decays after domination, where  $T_{\text{reh}}$  is given by Eq. (10). Combining the requirement that  $T_{\text{reh}} > T_{\text{BBN}}$  with the upper bound in Eq. (11) and using also Eq. (3), one finds the trivial bound:  $H_* > 10 T_{\text{BBN}}^2/m_P$ , which is way weaker than the bound in Eq. (8). The corresponding lower bound on  $g$  is

$$g > \sqrt{10} \left( \frac{m_P}{H_*} \right)^{1/2} \frac{T_{\text{BBN}}}{m_P}, \quad (15)$$

which is tighter than the lower bound in Eq. (11) if  $m < (T_{\text{BBN}}^2 m_P)^{1/3}$ . Finally, combining the condition  $\Gamma < H_{\text{dom}}$  with the lower bound in Eq. (11) one ends up again with the bound in Eq. (14).

In summary, the allowed range for the inflationary Hubble scale is

$$10 \text{ GeV} \leq H_* \leq 10^{14} \text{ GeV}, \quad (16)$$

while, for the decay coupling, the allowed range is

$$\max \left\{ \frac{T_{\text{BBN}}}{\sqrt{m_P m}}, \frac{m}{m_P} \right\} \lesssim g \lesssim \min \left\{ 1, \frac{m_P m}{T_{\text{BBN}}^2} \left( \frac{\sigma_*}{m_P} \right)^3 \right\} \quad (17)$$

Another set of bounds is due to the possible overproduction of gravitational waves (GWs) due to inflation. If inflation is followed by a phase whose equation of state is stiffer than radiation, the spectrum of relic gravitons features a spike (the slope grows with the frequency) for modes re-entering the horizon during the stiff phase [18]. Since in NO models there is a phase of kination right after inflation, high frequency gravitons re-entering the horizon during kination may disrupt BBN by increasing  $H$ . To avoid this, we require [13, 18]

$$I \equiv h^2 \int_{k_{\text{BBN}}}^{k_*} \Omega_{\text{GW}}(k) d \ln k \leq 2 \times 10^{-6}, \quad (18)$$

where  $\Omega_{\text{GW}}(k)$  is the GW density fraction with physical momentum  $k$  and  $h = 0.73$  is the Hubble constant  $H_0$  in units of 100 km/sec/Mpc. Using the spectrum  $\Omega_{\text{GW}}(k)$  as computed in [18], the above constraint can be written as [13]

$$I \simeq h^2 \alpha_{\text{GW}} \Omega_\gamma(k_0) \frac{1}{\pi^3} \frac{H_*^2}{m_P^2} \left( \frac{H_*}{H_{\text{dom}}} \right)^{2/3}, \quad (19)$$

where  $\alpha_{\text{GW}} \simeq 0.1$  is the GW generation efficiency during inflation and  $\Omega_\gamma(k_0) = 2.6 \times 10^{-5} h^{-2}$  is the density fraction of radiation at present on horizon scales.

If the curvaton field decays before domination, using  $H_{\text{dom}} \sim T_{\text{reh}}^2/m_P$  and Eq. (7), the bound in Eq. (18) becomes

$$\frac{H_*}{m} \left( \frac{\Gamma}{H_*} \right)^{1/3} \lesssim 24 \zeta^{-2} \sim 10^{10}. \quad (20)$$

If the curvaton decays after domination, using Eqs. (9) and (10), the bound in Eq. (18) becomes

$$\frac{H_*}{m_P} \left( \frac{H_*}{\Gamma} \right)^{1/3} \lesssim 1. \quad (21)$$

The above bounds may truncate further the ranges in Eqs. (16) and (17).

Let us quantify the above with a couple of specific examples. Firstly, let us choose  $H_* \sim 1$  TeV and  $m \sim 100$  GeV so that the bound in Eq. (3) is saturated. In this case Eq. (4) suggests that  $\sigma_* \sim 10^8$  GeV. Using this, Eq. (17) becomes  $10^{-13} \lesssim g \lesssim 10^{-4}$ . Using Eq. (9) we find  $H_{\text{dom}}/\Gamma \sim 10^{-20}/g^2$ , which means that the curvaton decays before domination if  $10^{-10} \lesssim g \lesssim 10^{-4}$ , whereas it decays after domination if  $10^{-13} \lesssim g \lesssim 10^{-10}$ . In the former case Eq. (7) gives  $T_{\text{reh}} \sim g^{-1/2} 10^{-5}$  GeV, while in the latter case Eq. (10) gives  $T_{\text{reh}} \sim g 10^{10}$  GeV. Hence, the allowed range for the reheating temperature is:  $1 \text{ MeV} < T_{\text{reh}} \lesssim 1 \text{ GeV}$ . Since  $\Gamma = g^2 m$ , it is straightforward to check that the bounds in Eqs. (20) and (21) are satisfied in the range shown in Eq. (17).

Now, let us choose  $\sigma_* \sim m_P$  and  $m \sim 100$  GeV, corresponding to a string axion as curvaton. In this case Eq. (4) suggests that  $H_* \sim 10^{14}$  GeV, which saturates the bound in Eq. (14) and corresponds to inflation at the grand unified scale. Using this, Eq. (17) becomes  $10^{-13} \lesssim g \lesssim 1$ . Then, Eq. (9) gives  $H_{\text{dom}}/\Gamma \sim g^{-2} \gtrsim 1$ . This means that the curvaton has to decay at or after domination. Thus, from Eq. (10) we have  $T_{\text{reh}} \sim g 10^{10}$  GeV, which results in the range:  $1 \text{ MeV} < T_{\text{reh}} \lesssim 10^{10}$  GeV. This time, however, the GW constraint turns out to be much tighter. Because the curvaton decays at or after domination, we need to use Eq. (21). With the chosen values, the bound can only be satisfied with  $g \sim 1$ , which saturates the upper bound in Eq. (11). Thus, the reheating temperature has to be  $T_{\text{reh}} \sim 10^{10}$  GeV, which seriously challenges gravitino constraints.

In conclusion, we have studied in detail the parameter space for the inflation scale in NO models with a minimal curvaton scenario being responsible for reheating the Universe as well as for the curvature perturbations. We have shown that low-scale inflation with  $H_*$  as low as 10 GeV is feasible in this case, regardless of whether the curvaton decays before or after dominating the Universe. Our results are independent of the particular form of the NO inflation model because, during the curvaton evolution, the inflaton sector is oblivious of the scalar potential, since it is dominated by the kinetic density of the inflaton.

### Acknowledgments

This work was supported (in part) by the European Union through the Marie Curie Research and Training Network "UniverseNet" (MRTN-CT-2006-035863) and by PPARC (PP/D000394/1).

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